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The Early Mirror Universe: Inflation, Baryogenesis, Nucleosynthesis and Dark Matter

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Abstract

There can exist a parallel 'mirror' world which has the same particle physics as the observable world and couples the latter only gravitationally. The nucleosynthesis bounds demand that the mirror sector should have a smaller temperature than the ordinary one. By this reason its evolution should be substantially deviated from the standard cosmology as far as the crucial epochs like baryogenesis, nucleosynthesis etc. are concerned. Starting from an inflationary scenario which could explain the different initial temperatures of the two sectors, we study the time history of the early mirror universe. In particular, we show that in the context of the GUT or electroweak baryogenesis scenarios, the baryon asymmetry in the mirror world should be larger than in the observable one and in fact the mirror baryons could provide the dominant dark matter component of the universe. In addition, analyzing the nucleosynthesis epoch, we show that the mirror helium abundance should be much larger than that of ordinary helium. The implications of the mirror baryons representing a kind of self-interacting dark matter for the large scale structure formation, the CMB anysotropy, the galactic halo structures, microlensing, etc. are briefly discussed.

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1 Introduction

The old idea that there can exist a hidden mirror sector of particles and interactions which is the exact duplicate of our visible world [1] has attracted a significant interest over last years, in particular being motivated by the problems of neutrino physics [2, 3], gravitational microlensing [4, 5], gamma ray bursts [6], ultra-high energy cosmic rays [7], flavour and CP violation [8], etc. The basic concept is to have a theory given by the product $G \times G'$ of two identical gauge factors with the identical particle contents, which could naturally emerge e.g. in the context of $E_8 \times E_8'$ superstring. Two sectors communicate through gravity and perhaps also via some other messengers. A discrete symmetry $P(G \leftrightarrow G')$ interchanging corresponding fields of G and G', so called mirror parity, implies that both particle sectors are described by the same Lagrangians.[1]

In particular, one can consider a minimal symmetry $G_{SM} \times G'_{SM}$ where $G_{SM} = SU(3) \times SU(2) \times U(1)$ stands for the standard model of observable particles: three families of quarks and leptons q_i, u_i^c, d_i^c; l_i, e_i^c $(i = 1, 2, 3)$ and the Higgs doublet ϕ, while $G'_{SM} = [SU(3) \times SU(2) \times U(1)]'$ is its mirror gauge counterpart with analogous particle content: fermions q_i', $u_i'^c$, $d_i'^c$; l_i', $e_i'^c$ and the Higgs ϕ'. (From now on all fields and quantities of the mirror (M) sector will have an apex to distinguish from the ones belonging to the observable or ordinary (O) world.) The mirror parity implies that all coupling constants (gauge, Yukawa, Higgs) have the same pattern in both sectors and thus their microphysics is the same.[2]

One could naively think that due to mirror parity the O- and M- particles should have the same cosmological densities, which would be in the immediate conflict with the Big Bang nucleosynthesis (BBN) bounds on the effective number of extra light neutrinos, $\Delta N_\nu < 1$ [11]: the mirror photons, electrons and neutrinos would give a contribution to the Hubble expansion rate equivalent to $\Delta N_\nu \simeq 6.14$. Therefore, the M-particle density in the early universe should be appropriately reduced. This situation is plausible if two following conditions are satisfied:

A. At the initial moment two systems are born with different densities. In particular, the inflationary reheating temperature in the M-sector should be lower than in the visible one, $T_R' < T_R$, which can be achieved in certain models [4, 7, 10].

B. The M- and O-particles interact very weakly, so that two systems do not come into the thermal equilibrium with each other in the early universe. This condition is automatically fulfilled if two worlds communicate only via the gravity. More generally, there could be other messengers like superheavy gauge singlet fields or light singlets of the moduli type. In either case, they should mediate the effective couplings between the O- and M- particles suppressed by a large mass factor $M \sim M_P$ or so.

If two sectors have different reheating temperatures, and if they do not come into the thermal contact at later stages, then during the universe expansion they evolve independently and approach the BBN epoch with different temperatures. Namely, the BBN bound $\Delta N_\nu < 1$ implies that $T'/T < 0.64$.

In this paper we study the comparative time history of two sectors in the early universe. We show that due to the temperature difference, in the mirror sector all key epochs as are the baryogenesis, nucleosynthesis, etc. proceed at somewhat different conditions than in the observable universe. The paper is organized as follows. In sect. 2 we describe an inflationary model which could provide different reheating temperatures between two sectors. The sect. 3 we show that the baryon asymmetry in the M-world generically should be larger than that in the O-sector either in the GUT or electroweak baryogenesis scenarios. Moreover, it is pretty plausible that M-baryons provide a significant fraction of the dark matter of the universe, i.e.

[1] In the brane world picture, the M-sector can be the same O-world realized on a parallel brane, $G' = G$ [9].

[2] The mirror parity could be spontaneously broken and the weak interaction scales $\langle\phi\rangle = v$ and $\langle\phi'\rangle = v'$ could be different, which leads to somewhat different particle physics in the mirror sector [3, 4, 7].

$\Omega'_B > \Omega_B$. In sect. 4 we study the chemical composition of the M-sector and find that it should be dominantly a helium world – the primordial abundance of the mirror ^{4}He should be $2-3$ times larger than of the ordinary one. In sect. 5 we briefly address the problem of the cosmological structure formation in the presence of M-baryons as a dominant dark matter component. Namely, M-baryons being a sort of self-interacting dissipative dark matter could provide interesting signatures on the CMB anisotropy, the large scale structure of the universe, the form of the galactic halos, microlensing, etc. Finally, in sect. 6 we summarize our findings.

2 Inflation and post-inflation

An attractive realization of the inflationary paradigm is provided by supersymmetric models of hybrid inflation. The symplest model is based on the superpotential $W = \lambda S(\Phi^2 - \mu^2)$ containing the inflaton field S and the additional 'orthogonal' field Φ, where λ is order 1 coupling constant and μ is a dimensional parameter of the order of the GUT scale [13]. The supersymmetric vacuum is located at $S = 0$, $\Phi = \mu$, while for the field values $\Phi = 0$, $S > \mu$ the tree level potential has a flat valley with an energy density $V = \lambda^2\mu^4$. Since the supersymmetry is broken by the non-vanishing F-term, $F_S = \lambda\mu^2$, the flat direction is lifted by radiative corrections and the potential of S gets a slope which is appropriate for the slow roll conditions. The COBE results on the CMB anisotropy imply that $V^{1/4} \simeq \epsilon^{1/4} \times 7 \cdot 10^{16}$ GeV, where $\epsilon \ll 1$ is a slow-roll parameter.

The genetic problem of the hybrid inflation models concerns the choice of the initial conditions [14, 15]. Namely, at the end of the Planck era the scalar S should have an initial value $\sim M_P$ while Φ must be zero with the high accuracy over a region much larger than the initial horizon size $\sim M_P^{-1}$. In other words, the initial field configuration should be located just on the bottom of the S valley, and so the energy density to start with should be $\sim \mu^4$, at least eight orders of magnitude smaller than the natural energy density M_P^4 at the Planck era. If the Φ field would have initial amplitude $\sim M_P$, then its oscillation should be damped within typical time $\sim M_P^{-1}$, otherwise it induces too a big curvature for the inflaton S and violate the slow-roll conditions, thus preventing the onset of the inflation.

Such a *Fine Tuning* can be avoided by a possible preheating before the onset of the inflation: the oscillating field Φ could decay into some light particles [15]. In this view, one can consider a model based on superpotential [15]:

$$W_{\rm preheat} = \lambda S(\Phi^2 - \mu^2) + g\Phi(\Psi^2 + \Psi'^2) \tag{1}$$

Now all fields can have initial values order M_P, and thus the initial energy density can be $\sim M_P^4$. At first instants, due to oscillation of Φ, the system behaves as matter dominated universe. Fast damping of Φ is a pretext of inflationary stage which allows the inflaton energy density $\sim \mu^4$ to dominate and then S can slowly roll to the origin. This function is carried by the second term in (1) – the oscillating orthogonal field Φ fastly decays into Ψ and Ψ' particles which have practically no contact to inflaton S. In addition, with vanishing Φ, also effective mass terms of Ψ's disappear and the latter fields start to behave as massless – they stop oscillating and freeze. In general, oscillations Ψ and Ψ' freeze at different amplitudes, typically $\sim \mu$, at which they are catched by the moment when their mass $g\Phi$ drops below the Hubble parameter (see Fig. 3 of ref. [15]). When slow roll ends up, all fields start to oscillate around their vacuum values, $S = 0$, $\Phi = \mu$, $\Psi, \Psi' = 0$, and reheat the universe.

Let us assume now that reheating occurs due to superpotential terms:

$$W_{\rm reheat} = h\Psi\phi_1\phi_2 + h\Psi'\phi'_1\phi'_2 \tag{2}$$

where $\phi_{1,2}$ and $\phi'_{1,2}$ are the Higgs doublet superfields respectively for the O- and M-sectors. Then different magnitudes of Ψ and Ψ' at the end of slow-roll phase should reflect into difference of reheating temperatures T_R and T'_R in two systems, simply because of the energy difference stored into Ψ and Ψ' oscillations which decay respectively into O- and M-Higgses. More detailed analysis of this mechanism will be presented elsewhere.[3]

We have also to make sure that after reheating two sectors do not come into the thermal equilibrium to each other, or in other words, that interactions between O- and M-particles are properly suppressed. In a supersymmetric theory this condition is fulfilled in a rather natural manner. In particular, the mixed terms of the lowest dimension in the superpotential are:

$$\frac{\beta_{ij}}{M}(l_i\phi)(l'_j\phi') + \frac{\beta}{M}(\phi_1\phi_2)(\phi'_1\phi'_2), \tag{3}$$

which are suppressed by a large mass factor $M \sim M_P$ or so, and thus are safe. The same holds true for soft supersymmetry breaking F- and D-terms like $[\frac{z}{M}(\phi_1\phi_2)(\phi'_1\phi'_2)]_F$, etc., where $z = m_S\theta^2$ is the supersymmetry breaking spurion with $m_S \sim 1$ TeV. As for the the kynetic mixing term $F^{\mu\nu}F'_{\mu\nu}$ between the field-strength tensors of the gauge factors $U(1)$ and $U(1)'$, it can be forbidden by embedding $G_{SM} \times G'_{SM}$ in the grand unified group like $SU(5) \times SU(5)'$.

Once the O- and M-systems are decoupled already after reheating, at later times t they will have different temperatures $T(t)$ and $T'(t)$, and so different energy and entropy densities:

$$\rho(t) = \frac{\pi^2}{30}g_*(T)T^4, \quad \rho'(t) = \frac{\pi^2}{30}g'_*(T')T'^4\ , \tag{4}$$

$$s(t) = \frac{2\pi^2}{45}g_s(T)T^3, \quad s'(t) = \frac{2\pi^2}{45}g'_s(T')T'^3\ . \tag{5}$$

The factors g_*, g_s and g'_*, g'_s accounting for the effective number of the degrees of freedom in two systems can in general be different from each other. During the universe expansion, the two sectors evolve with separately conserved entropies. Therefore, the ratio $x \equiv (s'/s)^{1/3}$ is time independent,[4] while the ratio of the temperatures in two sectors is simply given by:

$$\frac{T'(t)}{T(t)} = x \cdot \left[\frac{g_s(T)}{g'_s(T')}\right]^{1/3}\ . \tag{6}$$

The Hubble expansion rate is determined by the total energy density $\bar\rho = \rho + \rho'$, $H = \sqrt{(8\pi/3)G_N\bar\rho}$. Therefore, at a given time t in a radiation dominated epoch we have

$$H(t) = \frac{1}{2t} = 1.66\sqrt{\bar g_*(T)}\frac{T^2}{M_{Pl}} = 1.66\sqrt{\bar g'_*(T')}\frac{T'^2}{M_{Pl}} \tag{7}$$

in terms of O- and M-temperatures $T(t)$ and $T'(t)$, where

$$\bar g_*(T) = g_*(T)(1 + ax^4), \quad \bar g'_*(T') = g'_*(T')\left(1 + \frac{1}{ax^4}\right). \tag{8}$$

Here the factor $a(T, T') = [g'_*(T')/g_*(T)] \cdot [g_s(T)/g'_s(T')]^{4/3}$ takes into account that for $T' \neq T$ the relativistic particle contents of the two worlds can be different. However, except for very small values of x, we have $a \sim 1$. So hereafter we always take $\bar g_*(T) = g_*(T)(1 + x^4)$ and $\bar g'_*(T') = g'_*(T')(1 + x^{-4})$. In particular, in the modern universe we have $a(T_0, T'_0) = 1$, $g_s(T_0) =$

[3] For other scenarios of the asymmetric reheating, see [4, 7, 10].

[4] We assume that expansion goes adiabatically in both sectors and neglect the additional entropy production due to the possible weakly first order electroweak or QCD phase transitions.

$g'_s(T'_0) = 3.91$, and $x = T'_0/T_0$, where T_0, T'_0 are the present temperatures of the O- and M- relic photons.[5] In fact, x is the only free parameter in our model and it is constrained by the BBN bounds.

The observed abundances of light elements are in good agreement with the standard nucleosynthesis predictions, when at $T \sim 1$ MeV we have $g_* = 10.75$ as it is saturated by photons γ, electrons e and three neutrino species $\nu_{e,\mu,\tau}$. The contribution of mirror particles (γ', e' and $\nu'_{e,\mu,\tau}$) would change it to $\bar{g}_* = g_*(1+x^4)$. Deviations from $g_* = 10.75$ are usually parametrized in terms of the effective number of extra neutrino species, $\Delta g = \bar{g}_* - 10.75 = 1.75\Delta N_\nu$. Thus we have:

$$\Delta N_\nu = 6.14 \cdot x^4 \; . \tag{9}$$

In view of the present observational situation, a reliable bound is $\Delta N_\nu < 1$ [11], which translates as $x < 0.64$. This limit very weakly depends on ΔN_ν. E.g. $\Delta N_\nu < 1.5$ implies $x < 0.70$.

As far as $x^4 \ll 1$, in a relativistic epoch the Hubble expansion rate (7) is dominated by the O-matter density and the presence of M-sector practically does not affect the standard cosmology of the early ordinary universe. However, even if the two sectors have the same microphysics, the cosmology of the early mirror world can be very different from the standard one as far as the crucial epochs like baryogenesis, nuclesosynthesis, etc. are concerned. Any of these epochs is related to an instant when the rate of the relevant particle process $\Gamma(T)$, which is generically a function of the temperature, becomes equal to the Hubble expansion rate $H(T)$. Obviously, in the M-sector these events take place earlier than in the O-sector, and as a rule, the relevant processes in the former freeze out at larger temperatures than in the latter.

In the matter domination epoch the situation becomes different. In particular, we know that ordinary baryons can provide only a minor fraction of the present cosmological density, $\Omega_B = 0.01 - 0.06$, whereas the observational data indicate the presence of dark matter which amounts for $\Omega_m \sim 0.2 - 1$. It is interesting to question whether the missing matter density of the universe could be due to mirror baryons? In the next section we show that this situation can emerge in a pretty natural manner.

3 Baryogenesis and mirror baryon density

It is well known that a non-zero baryon asymmetry (BA) can be produced in the initially baryon symmetric universe if three following conditions are fulfilled: B-violation, CP violation and departure from the thermal equilibrium. Generally speaking, the baryogenesis scenarios can be divided in two categories in which the out of equilibrium conditions are provided (a) by the universe expansion itself, or (b) by fast phase transition and bubble nucleation. In particular, the latter concerns the electroweak baryogenesis schemes, while the former is typical for a GUT type baryogenesis or leptogenesis.

At present it is not possible to say which of the known mechanisms is responsible for the observed BA. We only know that the baryon to photon number density ratio $\eta = n_B/n_\gamma$ is restricted by the BBN constraints to the range $\eta = (2-6)\times 10^{-10}$. It is most likely that the BA in the M-world $\eta' = n'_B/n'_\gamma$ is produced by same mechanism and moreover, the rates of the B and CP violation processes are parametrically the same in both cases. However, the out of equilibrium conditions should be different since at relevant temperatures the universe expansion

[5] The frozen ratio of the neutrino and photon temperatures in the M-sector $r'_0 = T'_{\nu 0}/T'_0$ depends on the ν' decoupling temperature from the mirror plasma which scales approximatively as $T'_D \sim x^{-2/3} T_D$, where $T_D = 2-3$ MeV is the decoupling temperature of the usual neutrinos. Therefore, unless $x < 10^{-3}$, r'_0 has a standard value $r_0 = T_{\nu 0}/T_0 = (4/11)^{1/3}$. For $x < 10^{-3}$, T'_D becomes larger than the QCD scale $\Lambda \simeq 200$ MeV, so that due to the mirror gluons and light quarks u', d', s' contribution we would obtain $r'_0 = (4/53)^{1/3}$, $g'_s(T'_0) = 2.39$ and so $T'_0/T_0 \simeq 1.2x$. However, in the following such small values of x are not of our interest.

is faster for the M-sector. Below we show that by this reason η' typically emerges larger than η for either type (a) or (b) scenarios.

The M-baryons can be of the cosmological relevance if Ω'_B exceeds $\Omega_B = 3.67 \times 10^7 \eta h^{-2} = 0.01 - 0.06$, whereas $\Omega'_B > 1$ would overclose the universe. So we are interested in a situation when the ratio $\beta = \Omega'_B/\Omega_B$ falls in the range from 1 to about 100. Since $n'_\gamma = x^3 n_\gamma$, we obtain $\beta = x^3\eta'/\eta$. Therefore, $\eta' > \eta$ does not apriori mean that $\beta > 1$, and in fact there is a lower limit $x > 10^{-2}$ or so for the relevant parameter space. Indeed, it arises from $x^3 = \beta\eta/\eta'$ by recalling that $\eta \sim 10^{-9}$, while η' can be taken at most $\sim 10^{-3}$, the biggest value which can be principally realized in any baryogenesis scheme under the realistic assumptions.

3.1 GUT Baryogenesis

The GUT baryogenesis mechanism typically based on a superheavy boson X undergoing the B- and CP-violating decays into quarks and leptons. The following reaction rates are of relevance:
Decay: $\Gamma_D \sim \alpha_X M_X$ for $T \lesssim M_X$ or $\Gamma_D \sim \alpha_X M_X^2/T$ for $T \gtrsim M_X$, where α_X is the coupling strength of X to fermions and M_X is its mass;
Inverse decay: $\Gamma_{ID} \sim \Gamma_D (M_X/T)^{3/2} \exp(-M_X/T)$ for $T \lesssim M_X$ or $\Gamma_{ID} \sim \Gamma_D$ for $T \gtrsim M_X$;
The X boson mediated $2 \leftrightarrow 2$ processes: $\Gamma_S \sim n_X \sigma \sim A\alpha_X^2 T^5/(M_X^2 + T^2)^2$, where the factor A amounts for the possible reaction channels.

The final BA depends on a temperature at which X bosons go out from equilibrium. One can introduce a parameter which measures the effectiveness of the decay at the epoch $T \sim M_X$ [16]: $k = (\Gamma_D/2H)_{T=M_X} = 0.3\bar{g}_*^{-1/2}(\alpha_X M_{Pl}/M_X)$. The larger is k the longer equilibrium is maintained and the freeze-out abundance of X boson becomes smaller. Hence, the resulting baryon number to entropy ratio, $B = n_B/s \simeq 0.14\eta$ is a decreasing function of k. It is approximately given as $B \simeq \frac{\epsilon}{g_s} F(k, k_c)$, where ϵ is the CP violating factor and

$$F(k, k_c) = \begin{cases} 1 & \text{if } k < 1 \\ 0.3k^{-1}(\log k)^{-0.6} & \text{if } 1 < k < k_c \\ \sqrt{A\,\alpha_X\,k}\; e^{-\frac{4}{3}(A\,\alpha_X\,k)^{1/4}} & \text{if } k > k_c \end{cases} \qquad (10)$$

Here k_c is a critical value defined by equation $k_c(\log k_c)^{-2.4} = 300/(A\alpha_X)$. It distinguishes between the regimes $k < k_c$, in which inverse decay is relevant and $k > k_c$, in which instead $2 \leftrightarrow 2$ processes are the dominant reason for baryon damping.

In a general context, without referring to a particular model, it is difficult to decide which range of parameters k and k_c can be relevant for baryogenesis. One can impose only the most reasonable constraints $g_s(T = M_X) \geq 100$ and $\epsilon \leq 10^{-2}$, and thus $\epsilon/g_s < 10^{-4}$ or so. For a given mechanism responsible for the observed baryon asymmetry $B \sim 10^{-10}$, this translates into a lower bound $F(k, k_c) > 10^{-6}$.

The presence of the mirror sector practically does not alter the ordinary baryogenesis. The effective particle number is $\bar{g}_*(T) = g_*(T)(1 + x^4)$ and thus the contribution of M-particles to the Hubble constant at $T \sim M_X$ is suppressed by a factor x^4.

In the mirror sector everything should occur in a similar way, apart from the fact that now at $T' \sim M_X$ the Hubble constant is not dominated by the mirror species but by ordinary ones: $\bar{g}'_*(T') \simeq g'_*(T')(1 + x^{-4})$. As a consequence, we have $k' = (\Gamma'_D/2H)_{|T'=M_X} = kx^2$. Since the value of k_c is the same in the two sectors, the mirror baryon asymmetry can be simply obtained by replacing $k \to k' = kx^2$ in eq.10, i.e. $B' = n'_B/s' \simeq (\epsilon/g'_s)F(k' = kx^2, k_c)$. Since F is a decreasing function of k, then for $x < 1$ we have $F(kx^2, k_c) > F(k, k_c)$ and thus we conclude that the mirror world always gets a *larger* BA than the visible one, $B' > B$.

However, this does not apriori mean that Ω'_B is always larger than Ω_B. Since the entropy

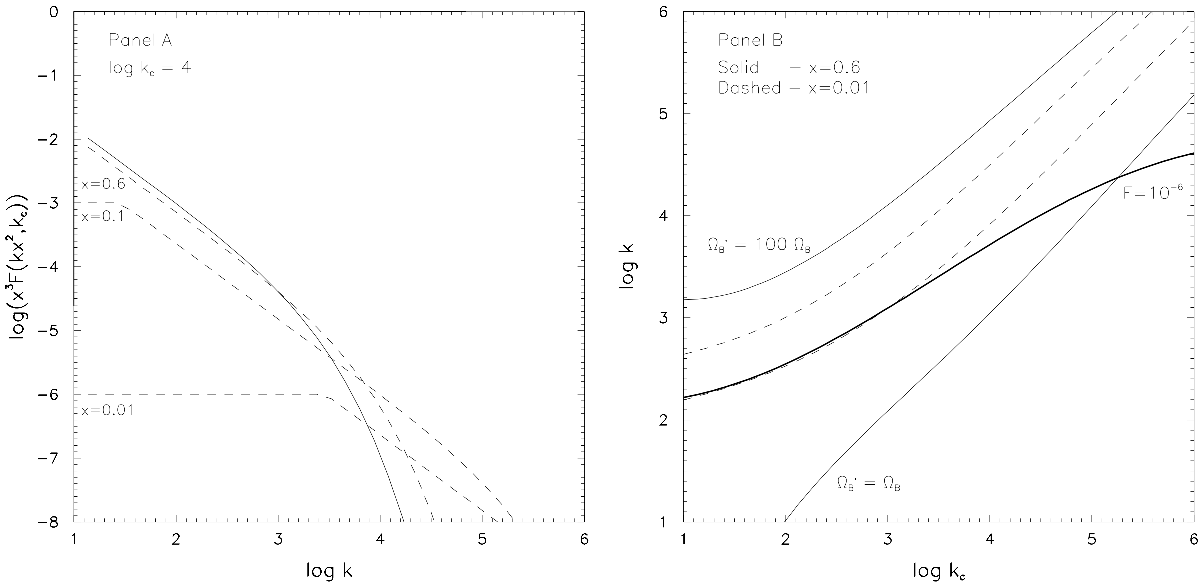

Figure 1: *Panel A*. The combination $x^3F(kx^2, k_c)$ as a function of k for $k_c = 10^4$ and $x = 0.6,\ 0.1,\ 0.01$ (dash). The solid curve corresponding to $x = 1$ in fact measures the possible BA in the ordinary world, $F(k, k_c) = (g_s/\epsilon)B(k)$. *Panel B*. The curves confining the parameter region in which $\beta = \Omega'_B/\Omega_B$ varies from 1 to 100, for $x = 0.6$ (solid) and for $x = 0.01$ (dash). The parameter area above thick solid curve corresponds to $F(k, k_c) < 10^{-6}$ and it is excluded by the observable value of η.

densities are related as $s'/s = x^3$, for the ratio $\beta = \Omega'_B/\Omega_B$ we have:

$$\beta(x) = \frac{n'_B}{n_B} = \frac{B's'}{Bs} = x^3\ \frac{F(kx^2, k_c)}{F(k, k_c)}\ . \tag{11}$$

The behaviour of the factor $x^3F(kx^2, k_c)$ as a function of k for different values of the parameter x is given in the Fig. 1A. Clearly, in order to have $\Omega'_B > \Omega_B$ the function $F(k, k_c)$ have to decrease faster than $k^{-3/2}$ between $k' = kx^2$ and k. Closer inspection of the function 10 reveals that the M-baryons can be overproduced only if k is order k_c or larger. In other words, the relevant interactions in the observable sector maintain equilibrium longer than in the mirror one, and thus ordinary BA can be suppressed by an exponential Boltzmann factor while the mirror BA could be produced still in non-exponential regime $k' < k_c$.

In Fig. 1B we show the parameter region in which $\beta = \Omega'_B/\Omega_B$ falls in the range $1 - 100$, in confront to the parameter area excluded by condition $F(k, k_c) > 10^{-6}$. We see that for $x = 0.6$ there is an allowed parameter space in which β can reach values up to 10, but $\beta = 100$ is excluded. For a limiting case $x = 10^{-2}$, as it was expected, the parameter space for $\beta > 1$ becomes incompatible with $F(k, k_c) > 10^{-6}$. For intermediate values of x, say $x \sim 0.1 - 0.3$, also the values $\beta \sim 100$ can be compatible.

The above considerations can be applied also in the context of leptogenesis. One should remark, however, that potentially both the GUT baryogenesis or leptogenesis scenarios are in conflict with the supersymmetric inflation scenarios, because of the upper limit on the reheating temperatures about $T_R < 10^9$ GeV from the thermal production of gravitinos [17]. Moreover, it was shown recently that the non-thermal gravitino production can impose much stronger limits, $T_R < 10^5$ GeV or so [18]. This problem can be fully avoided in the electroweak baryogenesis scenario, which instead is definitely favoured by the supersymmetry.

3.2 Electroweak Baryogenesis

The electroweak (EW) baryogenesis mechanism is based on the anomalous B-violating processes induced by the sphalerons which are quit rapid at high temperatures, but become much slower when temperature drops below 100 GeV. A succesfull scenario needs the first order EW phase transition and sufficient amount of CP violation, which conditions can be satisfied in the frames of the supersymmetric standard model, for certain parameter ranges [19].

The characteristic temperature scales of the electroweak phase transition are fixed entirely by the form of the finite temperature effective potential:

$$V(\phi, T) = D(T^2 - T_0^2)\phi^2 - ET\phi^3 + \lambda_T \phi^4 \tag{12}$$

where all parameters can be expressed in terms of the fundamental couplings in the Lagrangian. For large temperatures, $T \gg 100$ GeV, the electroweak symmetry is restored and $V(\phi, T)$ has a minimum at $\phi = 0$. With the universe expansion the temperature drops, approaching the specific values which define the sequence of the phase transition. These are all in the 100 GeV range and ordered as $T_1 > T_c > T_b > T_0$. Namely, below $T = T_1$ the potential gets a second local minimum $\phi_+(T)$. At the critical temperature $T = T_c$ the latter becomes degenerate with the false vacuum $\phi = 0$. At temperatures $T < T_c$ to the true vacuum state $\phi = \phi_+(T)$ becomes energetically favoured, and transition to this state can occur via thermal quantum tunneling, through the nucleation of the bubbles which then expand fastly, percolate and finally fill the whole space within a horizon by the true vacuum.

The bubble production starts when the free energy barrier separating the two minima becomes small enough. The bubble nucleation temperature T_b is defined as a temperature at which the probability for a single bubble to be nucleated within a horizon volume becomes order one:

$$P(T_b) = \omega \left(\frac{T_c}{H(T_c)}\right)^4 \left(1 - \frac{T_b}{T_c}\right) e^{-\frac{F_c(T_b)}{T_b}} \sim 1 \tag{13}$$

where $F_c(T)$ is the free energy and ω is an order one coefficient [19]. In particular, in the limit of thin wall approximation we have:

$$\frac{F_c(T)}{T} = \frac{64\pi}{81} \frac{E}{(2\lambda)^{3/2}} \left(\frac{T_c - T_0}{T_c - T}\right)^2 . \tag{14}$$

The condition (13) results into large values of $F_c(T_b)/T_b$, typically order 10^2.

Once the bubble nucleation and expansion rate is larger than the Hubble parameter, the out of equilibrium condition for anomalous B-violating processes is provided by the fast phase transition itself. The BA can be produced inside the bubbles due to the CP violation since the quarks and antiquarks have different reflection coefficients on the bubble wall. The baryogenesis rate is completely independent from the universe expansion and it occurs practically in one instant as compared to the cosmological time scale of this epoch.

As for the mirror sector, it is described by the same Lagrangian as the ordinary one and so the effective thermal potential of the mirror Higgs $V(\phi', T')$ has the same form as (12). Then the temperature scales which are defined entirely by the form of the effective potential should be exactly the same for O- and M-sectors. Namely, $T'_1 = T_1$ and $T'_c = T_c$.

The equation (13) is the same for the M-sector apart of the fact that the corresponding Hubble constant is different: $H(T' = T_c) = x^{-2} H(T = T_c)$. Therefore, we obtain:

$$\frac{F_c(T'_b)}{T'_b} = \frac{F_c(T_b)}{T_b} + 8 \log x \tag{15}$$

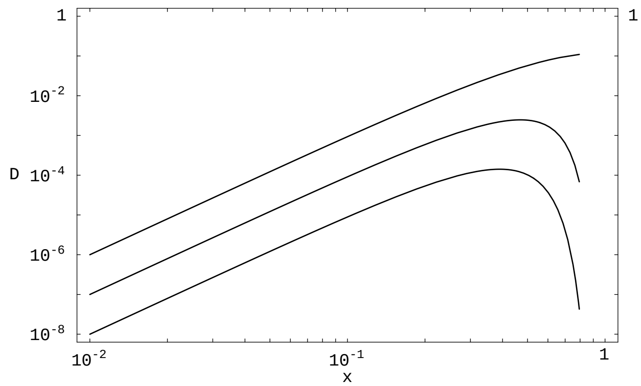

Figure 2: The contours of $\beta = \Omega'_B/\Omega_B$ in the plane of the parameters x and D, corresponding to $\beta = 1, 10$ and 100 from top to bottom.

which in turn tells that the bubble nucleation temperatures in two sectors are practically equal: $T'_b = T_b(1 + 0.01 \log x)$. (Clearly, the phase transition in M-sector occurs at earlier time than in O-sector: $t'_b \simeq x^{-2} t_b$.) The reason is that between the temperature scales T_c and T_0 the free energy $F_c(T)$ is a rapidly changing function but the change in temperature itself is insignificant. Hence, we expect that the initial BA's produced right at the phase transition to be the same in O- and M-sectors: $B(T = T_b) = B'(T' = T_b)$.

However, the instantly produced baryon number can be still washed out by the sphaleron interactions. The anomalous B-violation rate $\Gamma(T) \sim \exp[-F(T)/T]$, where $F(T)$ is the sphaleron free energy at finite T, may be large enough inside the bubble as far as the temperature is large. But it quickly falls as the temperature decreases, and baryon number freezes out as soon as $\Gamma(t)$ drops below $H(t)$ (7). The wash-out equation $dB/dt = -\Gamma(t)B$ can be rewritten as

$$\frac{dB}{B} = \frac{\Gamma(T)}{HT} dT \tag{16}$$

and integrated. Then the final BA in the O- and M-sectors can be expressed respectively as

$$B = B(T_b) D^{(1+x^4)^{-1/2}}, \qquad B' = B(T_b) D^{x^2(1+x^4)^{-1/2}}, \tag{17}$$

where $D < 1$ is the baryon number depletion factor:

$$D = \exp\left[-0.6 g_*^{-1/2} M_{Pl} \int_0^{T_b} dT \frac{\Gamma(T)}{T^3} \right] , \tag{18}$$

and $g_* \sim O(100)$ in the supersymmetric standard model. Thus we always have $B' > B$, while the M-baryon mass density relative to O-baryons reads:

$$\beta(x) = \frac{\Omega'_B}{\Omega_B} = x^3 \frac{B'}{B} = x^3 D^{-K(x)}, \qquad K(x) = \frac{1 - x^2}{\sqrt{1 + x^4}} . \tag{19}$$

Lacking a precise theory for non-perturbative sphaleron transitions in the broken phase, the exact value of D one cannot be calculated even in the context of concrete models. If $D \sim 1$, the wash-out is ineffective and practically all BA produced right at the bubble nucleation is conserved. In this case Ω'_B should smaller than Ω_B. However, if D is enough small, one can achieve sufficiently large Ω'_B. The contourplot for the parameters x and D for which β falls in the range $1 - 100$ is given in Fig. 2. For small x we have essentially $\beta \simeq x^3 D^{-1}$ and thus $100 > \beta > 1$ requires a depletion factor in the interval $D = (10^{-2} - 1)x^3$. Once again, for $x \sim 10^{-2}$ one needs the marginal values $D \sim 10^{-8} - 10^{-6}$ below which the observable BA $B \sim 10^{-10}$ cannot be produced at all.

4 Primordial Nucleosynthesis and mirror helium abundance

The time scales relevant for standard BBN are defined by the "freeze-out" temperature of weak interactions $T_W \simeq 0.8$ MeV ($t_W \sim 1$ s) and by the "deuterium bottleneck" temperature $T_N \simeq 0.07$ MeV ($t_N \sim 200$ s) [16]. When $T > T_W$, weak interactions transform neutrons into protons and viceversa and keep them in chemical equilibrium. The neutron abundance $X_n = n_n/n_B$, defined as the ratio of neutron to baryon densities, is given by $X_n(T) = [1 + \exp(\Delta m/T)]^{-1}$, where $\Delta m \simeq 1.29$ MeV is the neutron-proton mass difference. For $T < T_W$ the weak reaction rate $\Gamma_W \simeq G_F^2 T^5$ drops below the Hubble expansion rate $H(T) \simeq 5.5 T^2/M_{Pl}$, the neutron abundance freezes out at the equilibrium value $X_n(T_W)$ and it then evolves only due to the neutron decay: $X_n(t) = X_n(T_W)\exp(-t/\tau)$, where $\tau = 886.7$ s is the neutron lifetime.

At temperatures $T > T_N$, the process $p + n \leftrightarrow d + \gamma$ is faster than the universe expansion, and free nucleons and deuterium are in chemical equilibrium. The light element nucleosynthesis essentially begins when the system cools down to the temperature

$$T_N \simeq \frac{B_d}{-\ln(\eta) + 1.5\ln(m_N/T_N)} \simeq 0.07 \text{ MeV}, \tag{20}$$

where $B_d = 2.22$ MeV is the deuterium binding energy, and m_N is the nucleon mass. Below this temperature the deuterium abundance starts to grow which in turn allows to produce also the heavier nuclei. Nearly all neutrons present at this time are finally captured in ^{4}He nuclei, due to the large binding energy of the latter. Thus, the primordial ^{4}He mass fraction is:

$$Y_4 \simeq 2X_n(t_N) = \frac{2\exp(-t_N/\tau)}{1 + \exp(\Delta m/T_W)} \simeq 0.24. \tag{21}$$

As we have already discussed, the presence of the mirror sector with a temperature $T' \ll T$ has practically no impact the standard BBN. In fact, the limit $x < 0.64$ has been set by uncertainties of the present observational situation. In the mirror sector nucleosynthesis proceeds along the same lines. However, the impact of the O-world for the mirror BBN is dramatic.

For any given temperature T', now we have $H(T') \simeq 5.5(1 + x^{-2})T'^2/M_{Pl}$ for the Hubble expansion rate. Therefore, the freeze-out temperature $T'_W = (1 + x^{-4})^{1/6} T_W$ is larger than T_W, whereas the time scales as $t'_W = t_W/(1+x^{-4})^{5/6}$. In addition, η' is different from $\eta \simeq 5 \times 10^{-10}$. However, since T_N depends on baryon density only logarithmically (see (20)), the temperature T'_N remains essentially the same as T_N, while the time t'_N scales as $t'_N = t_N/(1+x^{-4})^{1/2}$. Thus, for the mirror ^{4}He mass fraction we obtain:

$$Y'_4 \simeq 2X'_n(t'_N) = \frac{2\exp[-t_N/\tau(1 + x^{-4})^{1/2}]}{1 + \exp[\Delta m/T_W(1 + x^{-4})^{1/6}]} \ . \tag{22}$$

We see that Y'_4 is an increasing function of x^{-1}. In particular, for $x \to 0$ one has $Y'_4 \to 1$.

In reality, eq. (22) is not valid for small x, since in this case deuterium production through reaction $n + p \leftrightarrow d + \gamma$ can become ineffective. By a simple calculation one can make sure that for $x < 0.3 \cdot (\eta' \times 10^{10})^{-1/2}$, the rate at which neutrons are captured to form the deuterium nuclei, $\Gamma'_N = n'_p \sigma_N \sim \eta' n'_\gamma \sigma_N$, where $n'_\gamma \sim T'^3$ is the M-photons density and $\sigma_N \simeq 4.5 \cdot 10^{-20}$ cm^3 s^{-1} is the thermal averaged cross section, becomes smaller than the Hubble rate $H(T')$ for temperatures $T' > T'_N$. In this case M-nucleosynthesis is inhibited, because the neutron capture processes becomes ineffective before deuterium abundance grows enough to initiate the synthesis of the heavier elements. Therefore, for any given η', Y'_4 first increases with increasing $1/x$, reaches a maximum and then starts to decrease. The true dependence of Y'_4 on the x computed for different values of η' by standard BBN code [20], is presented in Fig. 3. The

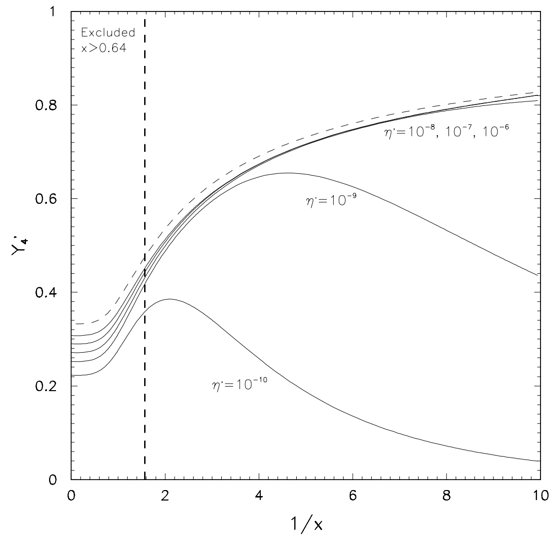

Figure 3: The primordial mirror ^{4}He mass fraction as a function of x. The dashed curve represents the approximate result of eq. (21). The solid curves obtained via exact numerical calculation correspond, from bottom to top, to η' varying from 10^{-10} to 10^{-6}.

Hubble expansion rate of the mirror world was implemented, for each value of x, by taking an effective number of extra neutrinos as $\Delta N_\nu = 6.14 \cdot x^{-4}$.

We have to remark, however, that in the most interesting situation when $\beta = \Omega'_B/\Omega_B = x^3\eta'/\eta > 1$, the condition $x < 0.3 \cdot (\eta' \times 10^{10})^{-1/2}$ is never fulfilled and the behaviour of Y'_4 is well described by the approximate formula (22). Hence, in this case Y'_4 is always bigger than Y_4. In other words, if dark matter of the universe is represented by the baryons of the mirror sector, it should contain considerably bigger fraction of primordial ^{4}He than the ordinary world.

5 Mirror baryons as dark matter

We have shown that mirror baryons could provide a significant contribution to the energy density of the universe and thus they could constitute a relevant component of dark matter. Immediate question arises: how the mirror baryon dark matter (MBDM) behaves and what are the differences from the more familiar dark matter candidates as the cold dark matter (CDM), the hot dark matter (HDM) etc. In this section we briefly address the possible observational consequences of such a cosmological scenario.

In the most general context, the present energy density contains relativistic (radiation) component Ω_r, non-relativistic (matter) component Ω_m and the vacuum energy density Ω_Λ (cosmological term). According the inflationary paradigm the universe should be almost flat, $\Omega_0 = \Omega_m + \Omega_r + \Omega_\Lambda \approx 1$, which well agrees with the recent results on the CMB anisotropy [21]. The Hubble parameter is known to be $H_0 = 100h$ km s^{-1} Mpc^{-1} with $h = 0.6 - 0.8$, and for redshifts of the cosmological relevance, $1 + z = T/T_0 \gg 1$, it becomes

$$H(z) = H_0 \left[\Omega_r(1+z)^4 + \Omega_m(1+z)^3\right]^{1/2} . \tag{23}$$

In the context of our model, the relativistic component $\Omega_r h^2 = 4.2 \times 10^{-5}(1 + x^4)$ is represented by ordinary photons and neutrinos and their mirror partners with a temperature scaled by a factor x. Contribution of the mirror species is negligible in view of the BBN constraint $x < 0.64$. As for the non-relativistic component, it should contain the O-baryon fraction Ω_B and the M-baryon fraction $\Omega'_B = \beta\Omega_B$. Although the other types of dark matter could also present,[6] here for simplicity we consider only the case $\Omega_m = \Omega'_B + \Omega_B$, when dark matter of the universe is entirely due to M-baryons.

At present it is not yet clear what is the matter fraction in the universe. Many observational data favour $\Omega_m \sim 0.3$ while the rest of the energy density is due to the cosmological term, $\Omega_\Lambda \sim 0.7$, but also $\Omega_m \sim 1$ cannot be excluded. For example, in our scenario one could have $\Omega_m = \Omega'_B + \Omega_B$ with $\Omega_B h^2 \simeq 0.02$ and $\Omega'_B h^2 \simeq 0.2$, for which just $\beta \simeq 10$ is needed.

The important moments for the structure formation are related to the matter-radiation equality (MRE) epoch which occurs at the redshift (we denote $(\Omega_m h^2)_{0.2} = (\Omega_m h^2/0.20)$):

$$1 + z_{\rm eq} = \frac{\Omega_m}{\Omega_r} \approx 2.4 \cdot 10^4 \frac{\Omega_m h^2}{1+x^4} = 4800 \times (\Omega_m h^2)_{0.2} \tag{24}$$

and to the plasma recombination and matter-radiation decoupling (MRD) epochs. The latter takes place only after the most of electrons and protons recombine into neutral hydrogen and the free electron number density n_e diminishes, so that the photon scattering rate $\Gamma_\gamma = n_e\sigma_T = X_e \eta n_\gamma \sigma_T$ drops below the Hubble expansion rate $H(T)$, where $\sigma_T = 6.65 \cdot 10^{-25}$ cm^2 is the Thomson cross section. In condition of chemical equilibrium, the fractional ionization $X_e = n_e/n_B$ is given by the Saha equation, which for $X_e \ll 1$ reads:

$$X_e \approx (1 - Y_4)^{1/2} \, \frac{0.51}{\eta^{1/2}} \left(\frac{T}{m_e}\right)^{-3/4} e^{-B/2T} \tag{25}$$

where $B = 13.6$ eV is the hydrogen binding energy. Thus we obtain the familiar result that in our universe the MRD takes place in the matter domination period, at the temperature $T_{\rm dec} \simeq 0.26$ eV which corresponds to redshift $1 + z_{\rm dec} = T_{\rm dec}/T_0 \simeq 1100$.

The MRD temperature in the M-sector $T'_{\rm dec}$ can be calculated following the same lines as in the ordinary one. Due to the fact that in either case the photon decoupling occurs when the exponential factor in eq. (25) becomes very small, we have $T'_{\rm dec} \simeq T_{\rm dec}$, up to small logarithmic corrections related to η', Y'_4 different from η, Y_4. Hence

$$1 + z'_{\rm dec} \simeq x^{-1}(1 + z_{\rm dec}) \simeq 1.1 \cdot 10^3 x^{-1} \tag{26}$$

so that the MRD in the M-sector occurs earlier than in the ordinary one. Moreover, for x less than $x_{\rm eq} = 0.23(\Omega_m h^2)^{-1}_{0.2}$, the mirror photons would decouple yet during the radiation dominated period (see Fig. 4).

Let us discuss now cosmological evolution of the MBDM. The relevant length scale for the gravitational instabilities is characterized by the mirror Jeans scale $\lambda'_J \simeq v'_s(\pi/G\rho)^{1/2}$, where $\rho(z)$ is the matter density at a given redshift z and $v'_s(z)$ is the sound speed in the M-plasma:

$$v'_s(z) \simeq \frac{c}{\sqrt{3}} \sqrt{1 + \frac{3}{4}\frac{\rho'_B}{\rho'_\gamma}} \approx \frac{c}{\sqrt{3}} \times \left[1 + \frac{3(1+z'_s)}{4x^4(1+z)}\right]^{-1/2} . \tag{27}$$

Here $1 + z'_s$ stands for a redshift at which the M-matter density $\rho'_B(z) \propto \Omega_m(1+z)^3$ becomes equal to the M-radiation density $\rho'_\gamma(z) \propto \Omega'_\gamma(1+z)^4 \simeq 2.5 \times 10^{-5}(1+z)^4 x^4$:

$$1 + z'_s = \frac{\Omega'_B}{\Omega'_\gamma} \approx 4.0 \cdot 10^4 \frac{\Omega_m h^2}{x^4} = 8000 \times x^{-4} (\Omega_m h^2)_{0.2} \tag{28}$$

[6]For example, some HDM fraction would emerge if any of the observable neutrinos has a mass in the eV range. In the context of supersymmetry, the CDM component could be provided by the lightest supersymmetric particle (LSP) from the O-sector. It is interesting to remark that mass fractions of their mirror partners would scale respectively as $\Omega'_\nu = x^3\Omega_\nu$ and $\Omega'_{\rm LSP} = x\Omega_{\rm LSP}$, and so their contribution should be irrelevant in either case.

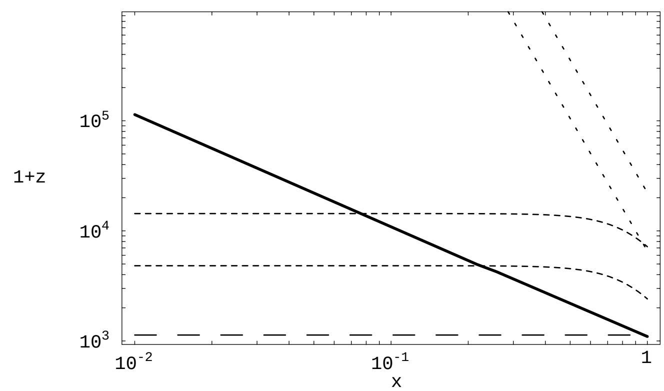

Figure 4: The M-photon decoupling redshift $1+z'_{dec}$ as a function of x (solid). The long-dash line marks the ordinary MRD redshift $1+z_{\rm dec}=1100$. We also show the MRE redshift $1+z_{\rm eq}$ and the mirror sound speed break redshift $1+z'_s$ for $\Omega_m h^2=0.2$ (respectively lower dash and lower dott) and $\Omega_m h^2=0.6$ (upper dash and dott).

(see Fig. 4). Since the M-plasma contains more baryons and less photons than the ordinary one, for the redshifts of the cosmological importance around $z_{\rm eq}<z'_s$ we have $v'_s \ll c/\sqrt{3}$, quite in contrast with the ordinary world, where $v_s \simeq c/\sqrt{3}$ practically till the photon decoupling.

The M-baryon Jeans mass $M'_J=\frac{\pi}{6}\rho_m\lambda_J'^3$ reaches the maximal value at $z=z'_{\rm dec}\simeq 1100/x$, $M'_J(z'_{dec})\simeq 6{\cdot}10^{17}{\times}x^6[1{+}(x_{\rm eq}/x)]^{-3/2}(\Omega_m h^2)^{-2}_{0.2}\,M_\odot$. Notice, however, that M'_J becomes smaller than the Hubble horizon mass $M_H=\frac{\pi}{6}\rho H^{-3}$ starting from a redshift $z_c=750x^{-4}(\Omega_m h^2)_{0.2}$. For $x=0.64$ it is about $z_{\rm eq}$ but it sharply increases for smaller values of x. So, the density perturbation scales which enter horizon about $z\sim z_{\rm eq}$ have mass larger than M'_J and thus undergo uninterrupted linear growth immediately after $t=t_{\rm eq}$. The smaller scales for which $M'_J>M_H$ instead would first oscillate. Therefore, the large scale structure formation is not delayed even if the mirror MRD epoch did not occur yet, i.e. even if $x>x_{\rm eq}$. The density fluctuations start to grow in the M-matter and the visible baryons are involved later, after being recombined, when they rewritte the spectrum of already developed mirror structures.

The main feature of the MBDM scenario is that the M-baryon density fluctuations should undergo the strong collisional damping around the time of M-recombination. The photon diffusion from the overdense to underdense regions induce a dragging of charged particles and wash out the perturbations at scales smaller than the mirror Silk scale $\lambda'_S\simeq 10\times f(x)(\Omega_m h^2)^{-3/4}_{0.2}$ Mpc, where $f(x)=x^{5/4}$ for $x>x_{\rm eq}$, and $f(x)=(x/x_{\rm eq})^{3/2}x_{\rm eq}^{5/4}$ for $x<x_{\rm eq}$.

Thus, the density perturbation scales which can run the linear growth after the MRE epoch are limited by the length λ'_S. This could help in avoiding the excess of small scales (of few Mpc) in the large scale power spectrum without tilting the spectral index.[7] The smallest perturbations that survive the Silk damping will have the mass $M'_S\sim x^{15/4}(\Omega_m h^2)^{-5/4}_{0.2}10^{13}\,M_\odot$, which should be less than $2\times10^{12}\,M_\odot$ in view of the BBN bound $x<0.64$. Interestingly, for $x\sim x_{\rm eq}$ we have $M'_S\sim 4\times10^{10}\,(\Omega_m h^2)^{-5}_{0.2}\,M_\odot$, a typical galaxy mass.

To some extend, the cutoff effect is analogous to the free streaming damping in the case of warm dark matter (WDM), but there are important differences. The point is that alike usual baryons, the MBDM should show acoustic oscillations whith an impact on the large scale power spectrum. In particular, it is tempting to imagine that the M-baryon oscillation effects are related to the anomalous features observed at $100h^{-1}$ Mpc clustering [22].

[7]One should keep on mind also the possibility of delaying the MRE moment itself, e.g. by taking $\Omega_m h^2\sim 0.2$. This would correspond to the case when M-baryons provide only a fraction of the present energy density, say $\Omega_m\sim 0.3$, while the rest is due to the cosmological term, $\Omega_\Lambda\sim 0.7$.

In addition, the MBDM oscillations transmitted via gravity to the ordinary baryons, could cause observable anomalies in the CMB angular power spectrum for l's larger than 200. This effect can be observed only if the M-baryon Jeans scale λ'_J is larger than the Silk scale of ordinary baryons, $\lambda_S \simeq 3\beta^{1/2}(\Omega_m h^2)^{-3/4}_{1/5}$ Mpc, which sets a principal cutoff for CMB oscillations around $l \sim 1200$. As we have seen above, this would require enough large values of x, near the upper bound set by the BBN constraints: $x \simeq 0.6$ or so. The detailed analysis of this effect will be given elsewhere. In our opinion, together with the possible effects on the large scale power spectrum, it can provide a most direct test of the MBDM scenario and can be verified by the next CMB experiments with higher sensitivity.

Clearly, for small x the M-matter recombines much before the MRE moment, and thus it should rather manifest as the CDM as far as the large scale structure is concerned. However, there still can be crucial difference at smaller scales which already went non-linear, like galaxies.

Since in our scenario dark matter in galaxies and clusters are mirror baryons, one can question whether their distribution in halos is different from that of the CDM? Namely, simulations show that the CDM forms triaxial halos with a density profile too clumped towards the center, and overproduce the small substructures within the halo. As for the MBDM, it constitutes a sort of collisional dark matter and thus potentially could avoide these problems, at least clearly the one related to the excess of small substructures.

The halo distribution in galaxies depends on the mass M and on the self-scattering cross section σ of dark matter. In our case it mainly consists of the mirror hydrogen atoms, and so $M_H \simeq 1$ GeV and $\sigma_H \sim 10^{-16}$ cm^2. At the first glance, this is in strong discrepancy with the range $\sigma/M \sim 10^{-23} - 10^{-24}$ cm^2/GeV preferred by the analysis of ref. [24].[8]

However, one has to take into account the possibility that during the galaxy evolution the bulk of the M-baryons could fastly fragment into the stars and only part of them could be in gaseous state, which in addition can be ionized. A difficult question which we cannot answer here is related to the stellar formation in the M-sector, also taking into account that its temperature/density conditions and chemical contents are much different from the ordinary ones. In any case, the fast stellar formation would extint the mirror gas and the protogalaxy (which at certain moment essentially becomes the collisionless system of the mirror stars) could develope a typical elliptical structure.[9] Certainly, in this consideration also the galaxy merging process should be taken into account.

However, the inner part of halo can still contain the large amount of the ionized mirror gas and it is not excluded that it can have a quasi-spherical form. Even if the stellar formation is very efficient, the massive mirror stars in the dense central region fastly evolve[10] and explode as supernovae, leaving behind the compact objects like neutron stars or black holes,[11] and reproducing the mirror gas also including the heavier elements. Although the cross section σ_H is large, it does not necessarily implies that the galaxy core will collapse within a dynamical time, since the inner halo should be opaque for M-particles. They undergo many scatterings and escape from the system via diffusion, so the energy drain can be small enough and the instability time can substantially exceed the age of the universe [25].

In the galactic halo (which in fact is the mirror galaxy) the mirror stars should be observed as Machos in gravitational microlensing [4, 5]. Leaving aside the difficult question of the initial

[8] In the case of asymmetric M-world with $\zeta = v'/v \gg 1$, the mirror electron mass should scale as $m'_e \simeq \zeta m_e$ while the nucleon mass remains ~ 1 GeV. Therefore, the elastic scattering cross-section of the M-hydrogen atoms scales as $\sigma'_H \sim \zeta^{-2}\sigma_H$ and so for $\zeta \ll 1$ one could obtain a reasonably small cross section.

[9] As for the O-matter, within the dark M-matter halo it should typically show up as an observable elliptic or spiral galaxy, but some anomalous cases can be also possible.

[10] Since mirror matter contains more helium, the mirror stars may evolve faster than the ordinary ones.

[11] Another tempting issue is whether the M-matter itself could help in producing big central black holes, with masses $\sim 10^7\ M_\odot$, which are thought to be main engines of the active galactic nuclei.

stellar mass function, one can remark that once the mirror stars could be very old and evolve faster than the ordinary ones, it is suggestive to think that most of massive ones, with mass above the Chandrasekhar limit $M_{\rm Ch} \simeq 1.5\ M_\odot$ have already ended up as supernovae, so that only the lighter ones remain as the microlensing objects.[12] The recent data indicate the average mass of Machos around $M \simeq 0.5\ M_\odot$, which is difficult to explain in terms of the brown dwarves with masses below the hydrogen ignition limit $M < 0.1 M_\odot$ or other baryonic objects [23]. Perhaps, this is the first observational evidence of the mirror matter?

It is also plausible that in the galactic halo some fraction of mirror stars exists in the form of compact substructures like globular or open clusters. In this case, for a significant statistics, one could observe interesting time and angular correlations between the microlensing events.

6 Conclusions

We have discussed cosmological implications of the parallel mirror world with the same microphysics as the ordinary one, but having smaller temperature, with the limit $T' < 0.64T$ set by the BBN constraints. Therefore, the M-sector contains less relativistic matter (photons and neutrinos) than the O-sector, $\Omega'_r \ll \Omega_r$, and so in the relativistic expansion epoch the cosmological energy density is dominantly due the ordinary component, while the mirror one gives a negligible contribution. We have shown that by this reason the mirror sector should produce larger baryon asymmetry than the observable one, in the context of the GUT or electroweak baryogenesis scenarios. This can converge in intriguing complementary situation for the non-relativistic epoch: the mirror baryons can have a bigger density than the observable ones, $\Omega'_B \gg \Omega_B$, so that they can constitute a dominant dark matter of the universe.

We have also shown that the BBN epoch in the mirror world proceeds differently from the ordinary one, and it predicts the mirror helium abundance in the range $Y'_4 = 0.4 - 0.8$, considerably larger than the observable $Y_4 \simeq 0.24$.

Unfortunately, we cannot exchange the information with the mirror physicists and combine our observations. (Afterall, since two worlds have the same microphysics, the life should be possible also in the mirror sector.) However, there can be many possibilities to disentangle the cosmological scenario of two parallel worlds with the the future high precision data concerning the large scale structure, CMB anisotropy, structure of the galaxy halos, gravitational microlensing, oscillation of the neutrinos or other neutral particles into their mirror partners, etc.

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[12] The M-supernovae explosion in our galaxy cannot be directly seen by ordinary observer, however it could be observed in terms of gravitational waves. In addition, if the M- and O-neutrinos are mixed [2, 3], it can lead the observable neutrino signal, which could be also accompanied by the weak gamma ray burst [6].

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